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Probing chiral dynamics by charged-pion correlations

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The environment generated in the midrapidity region of a high-energy nuclear collision endows the pionic degrees of freedom with a time-dependent effective mass. Its specific evolution provides a mechanism for the production of back-to-back charge-conjugate pairs of soft pions which may present an observable signal of the nonequilibrium dynamics of the chiral order parameter.

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High-energy nuclear collisions are expected to produce transient systems within which chiral symmetry is approximately restored and the matter is partially deconfined. The identification and exploration of such a novel phase of matter is currently a major experimental goal and the efforts have intensified with the recent commissioning of the Relativistic Heavy Ion Collider (RHIC) at BNL.

Through the past decade, it has been speculated that the rapid expansion of the collision zone, after an approximate restoration of chiral symmetry has occurred, may produce long-wavelength isospin-polarized agitations of the pionic field, commonly referred to as disoriented chiral condensates (DCC), which in turn should lead to anomalies in the resulting pion multiplicity distribution. Reviews of this topic are given in Refs. [1–3].

Initially the focus was on the expected broadening in the distribution of the neutral pion fraction (the number of neutral pions divided by the total pion number) but experimental efforts mounted to search for such a signal yielded a null result [4,5]. It has meanwhile become increasingly clear from model calculations that any DCC effect is carried exclusively by relatively soft pions and thus the signal can be significantly enhanced by limiting the analysis to those pions. Unfortunately, the neutral pions decay into photon pairs and it is practically impossible to perform a reconstruction that would permit the determination of the individual speeds. As a consequence, the neutral pion fraction is less than ideal as an experimental indicator.

Nevertheless, the neutral pion fraction has received renewed attention by the recent realization [6] that it may be enhanced by the axial anomaly in a manner that depends in a characteristic way on the angle of emission with respect to the scattering plane. However, it is still an open question to what degree this interesting signal can in fact be extracted by suitable reconstruction of the photon data.

Furtunately, recent detailed dynamical calculations suggest that there may be suitable signals that are based exclusively on the charged pions and which would therefore better lend themselves to experimental detection. In particular, simulations within the semiclassical linear σ model [7] for cylindrical sources endowed with a longitudinal Bjorken scaling expansion identified a number of simpler candidate observables. One is the transverse spectral profile which is expected to exhibit a marked enhancement below $\approx 300\ \mathrm{MeV}/c$ due to the parametric amplification by the time modulation of the mass. However, an experimental extraction of such an effect is hampered by the contributions from a number of additional sources of soft pions. Thus, any observed signal may be less specific and would need to be correlated with other signatures. A more promising candidate observable identified in Ref. [7] is the fluctuation in the multiplicity of soft pions which was found to exhibit a significant anomalous increase with the order of the correlation. Since a subsequent study [8] has shown that such an effect is absent in both RQMD and HIJING (both widely used event generators using only conventional physics input), it appears that this observable has a larger specificity and thus it may be worthwhile to pursue it in the data analysis.

To the above possibilities for DCC signals, the present communication adds a novel suggestion for an observable that may be particularly suitable as an indicator of the phenomenon. It is based on the key feature that the basic DCC production mechanism creates neither momentum nor charge and so a signature may exist in the form of a large-angle correlation between oppositely charged soft pions. We first describe the basic effect, then illustrate the signal by suitable model calculations, and finally discuss the prospects for its observability.

The initial violent collision of the two approaching nuclei is followed by a rapid expansion, which at first proceeds in the longitudinal direction and then gradually builds up transversely as well. The pionic degrees of freedom then experience an environment that is changing accordingly. Generally, the effect of the environment can be approximately accounted for by an in-medium effective mass which depends both on the degree of agitation and on the chiral order parameter. Since the system is steadily cooling down while the order parameter reverts from its initial small value to its large vacuum value in a nonequilibrium fashion, the effective pion mass has then a correspondingly intricate evolution, displaying an overall decay towards the free mass overlaid by the effect of the oscillations by the relaxing order parameter [7,9].

Considerable insight into this coupled evolution can be gained from numerical studies with the linear σ model and though the details depend on the specific treatment, the emerging qualitative picture is fairly robust. It is therefore of interest to consider what implications this type of time evolution of the medium-modified mass may have on the pion observables. Scenarios with a time-dependent effective mass occur in many areas of physics, for example, in connection with the inflationary universe, and they can be treated to various degrees of refinement. For our present purposes it suffices to recognize a few general properties of such processes.

In order to bring these out most clearly, we consider first the simple case where the environment, and hence the effective mass, is spatially uniform, as is approximately the case in the interior of the collision zone. It is then obvious that although the time dependence of the mass may generate considerable agitation, this agency cannot add any net momentum. Thus any pions produced by the mechanism must be formed pairwise and moving in opposite directions. Furthermore, by a similar reasoning, the time dependence of the mass does not add any charge, so the produced pairs must be oppositely charged. Thus, the particles generated by an arbitrary time dependence in a uniform medium are charge-conjugate back-to-back pairs. This basic feature may be exploited as a probe of the chiral dynamics.

In order to illustrate the effect, we stay first with the spatially uniform scenario which is the simplest because the Hamiltonian decouples so that each degenerate pair of oppositely moving modes can be considered separately. A quantum-field treatment is then relatively simple [10]. The eigenmodes for any given value of the effective mass, μ, can be obtained by a suitable squeezing of the standard free modes. Since these modes (which we shall denote as quasiparticles) diagonalize the problem for the particular value of μ they provide a particularly instructive basis for the analysis of the evolution. We use $\hat{A}_{\mathbf{k}}$ to denote the annihilation operator for the quasiparticle modes corresponding to the initial value of the effective mass, μ_i. The Heisenberg representation of the corresponding time-dependent quasiparticle mode is then

$$\hat{A}^{\nu}_{\mathbf{k}}(t)=U_k(t)\hat{A}^{\nu}_{\mathbf{k}}+V_k(t)^*(\hat{A}^{\bar{\nu}}_{-\mathbf{k}})^{\dagger}. \tag{1}$$

Here the subscript **k** denotes the momentum and the superscript $\nu=\pm$ denotes the charge state (with $\bar{\nu}\equiv-\nu$). The time dependence of the quasiparticle operators $\hat{A}^{\nu}_{\mathbf{k}}(t)$ derive both from the usual Heisenberg evolution and from the continual redefinition of the quasiparticle basis. The Bogoliubov coefficients $U_k(t)$ and $V_k(t)$ are given in terms of the corresponding mode functions which in turn are determined by the equation of motion for the pion field. Being dependent only on the mode frequency ω_k, they are the same for all charge components ν and for all momenta with the same magnitude k. The time evolution mixes particle states of forward-going positive pions with hole states of backward-going negative pions, a reflection of the key conservation properties.

Let us now for simplicity assume that the system can initially be described as a thermal ensemble. The corresponding occupancy of a given initial quasiparticle mode is then $N_k=1/[\exp(\omega_k/T)+1]$, where T is the specified temperature. The expected occupancy of the mode at a later time t readily follows from Eq. (1):

$$\begin{aligned}N^{\nu}_{\mathbf{k}}(t)&\equiv\langle\hat{N}^{\nu}_{\mathbf{k}}(t)\rangle\\&\equiv\langle\hat{A}^{\nu}_{\mathbf{k}}(t)^{\dagger}\hat{A}^{\nu}_{\mathbf{k}}(t)\rangle\\&=|U_k(t)|^2N_k+|V_k(t)|^2(N_k+1).\end{aligned} \tag{2}$$

With a bit more elementary operator algebra, it is possible to obtain the following two equivalent expressions for the quasiparticle correlation coefficient:

$$\begin{aligned}\sigma^{\nu'\nu}_{\mathbf{k}'\mathbf{k}}(t)&\equiv\langle\hat{N}^{\nu'}_{\mathbf{k}'}(t)\hat{N}^{\nu}_{\mathbf{k}}(t)\rangle-N^{\nu'}_{\mathbf{k}'}(t)N^{\nu}_{\mathbf{k}}(t)\\&=N_k\bar{N}_k\delta_{\nu'\nu}\delta_{\mathbf{k}'\mathbf{k}}+|U_k(t)|^2|V_k(t)|^2\\&\quad\times(2N_k+1)^2[\delta_{\nu'\nu}\delta_{\mathbf{k}'\mathbf{k}}+\delta_{\nu'\bar{\nu}}\delta_{\mathbf{k}',-\mathbf{k}}],\end{aligned} \tag{3}$$

$$\begin{aligned}&=N_{\mathbf{k}}(t)\bar{N}_{\mathbf{k}}(t)\delta_{\nu'\nu}\delta_{\mathbf{k}'\mathbf{k}}+|U_k(t)|^2|V_k(t)|^2\\&\quad\times(2N_k+1)^2\delta_{\nu'\bar{\nu}}\delta_{\mathbf{k}',-\mathbf{k}}.\end{aligned} \tag{4}$$

In the first relation, the term $N\bar{N}$ is the initial Bose-Einstein autocorrelation for the occupancy of the given mode while the second term is the additional correlation introduced by the time dependence of the effective mass, $\mu(t)$, and its specific form reflects the fact that the autocorrelation is increased in concert with the back-to-back correlation between charge-conjugate pairs. The last relation shows that the resulting autocorrelation still has the familiar Bose-Einstein form, $N(t)\bar{N}(t)$.

It may be instructive at this point to discuss the present effect in relation to the medium-generated correlations studied by Asakawa, Csörgö, and Gyulassy [11] in a different context. In both scenarios, the starting point is a medium-modified effective pion mass which, upon diagonalization of the Hamiltonian, causes the quasiparticle eigenmodes to be superpositions of back-to-back correlated pairs of charge-conjugate free pions. If the system is assumed to disassemble very suddenly, as is being assumed in Ref. [11], this initial state remains unchanged and the result is the appearance of free pions endowed with the corresponding correlations. In the opposite extreme, when the reversion to the free scenario is adiabatic, all occupancy numbers (with respect to the evolving quasiparticle basis) remain constant in time as these modes gradually lose their composite character [12]. The result is then a gas of free pions with a spectral distribution reflecting the initial temperature of the in-medium quasiparticle gas. The present study involves the general dynamical scenario lying in between these two idealized extremes. The temporal evolution of the medium can thus cause real excitations to occur, thereby increasing the number of quasiparticle quanta present. The result of this complicated evolution is then a gas of free pions displaying a residue of the initial in-medium quasiparticle correlations (those considered in Ref. [11]), augmented by the pion pairs produced by parametric amplification.

These features are illustrated in Fig. 1. In order to emulate a uniform, longitudinally expanding environment, we use a simple time dependence of the effective mass, $\mu^2(t)=m^2+(\mu_i^2-m^2)\exp(-t/t_0)[1+\cos(\omega_{\rm osc}t)]/2$. With $\mu_i=2m$, t_0

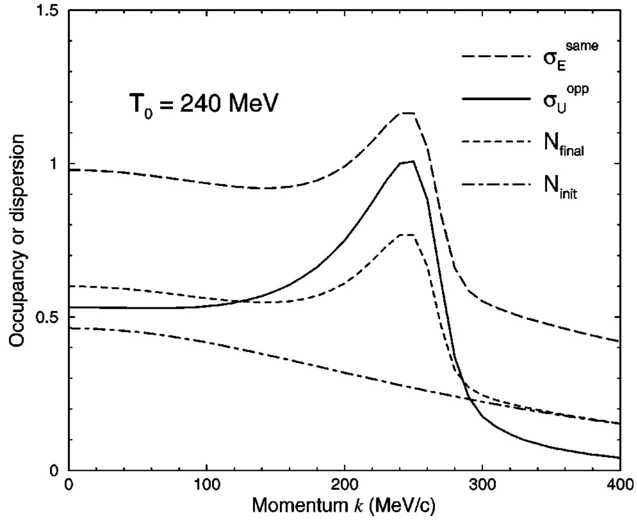

FIG. 1. The amplification and correlation obtained for an idealized spatially constant mass function $\mu(t)$ emulating an expanding scenario that starts from an environment in which the order parameter is $\sigma_0 \approx 26$ MeV (corresponding to $T_0 \approx 240$ MeV). Further explanation is given in the text.

$=5$ fm/c, and $\omega_{osc}=m_\sigma$, this expression approximates the evolution from an initial environment in which the order parameter has the value $\sigma_0 \equiv \langle \overline{q}q \rangle \approx 26$ MeV (corresponding to a temperature of $T_0 \approx 240$ MeV in the semiclassical linear σ model [7]).

Figure 1 shows how the initial Bose-Einstein quasiparticle occupation numbers N_{init} (dot-dashed curve) are being enhanced at momenta below ≈ 300 MeV/c (short dashes). The resulting peak structure in N_{final} is a reflection of the regularity of the order parameter oscillations: particular amplification is experienced by the modes that have their frequency ω_k in the neighborhood of half the σ mass, corresponding to $k \approx 260$ MeV for $m_\sigma = 600$ MeV. The autocorrelation in the mode occupancies (long dashes) exhibits a similar structure (since it is given by $N\overline{N}$). It is noteworthy to mention that although the correlation between oppositely moving charge-conjugate pion pairs (solid curve) is generally smaller than the autocorrelation, as Eq. (3) dictates, it acquires a comparable magnitude in the region of strongest amplification (where the dynamically produced pions dominate over those originally present). On the other hand, this correlation quickly subsides as the momentum moves above ≈ 300 MeV/c, bringing out the fact that the effect is confined to the soft regime.

The above discussion applies to the idealized scenario of an entirely uniform environment (implying, for example, that the singles yield is isotropic). It would be more realistic to consider a mass function that deviates from the free value only within a finite volume representing the agitated region. When the environment has a spatial dependence, the pions experience forces which tend to erode the clear back-to-back correlation pattern. Thus it is important to ascertain the importance of this effect.

It was recently shown how the quantum-field treatment of Ref. [10] can be extended to nonuniform scenarios as well, still within the effective mass approximation [12]. In that work, an illustrative application was made to a simple one-dimensional scenario, in which the effective mass function has a diffuse (Woods-Saxon–like) spatial profile with an overall temporal modulation that approximates the effective mass function obtained from dynamical simulations of Bjorken rods as described in Ref. [7]. We have employed that approach, for the same simplified scenario, to obtain the time evolution of the two-body correlation function and the result is illustrated in Fig. 2. As expected, the correlation pattern seen in Fig. 1 has become somewhat eroded, due to the spatial dependence of the mass function, but the basic structure largely persists. Thus there is reason to hope that it may be experimentally observable.

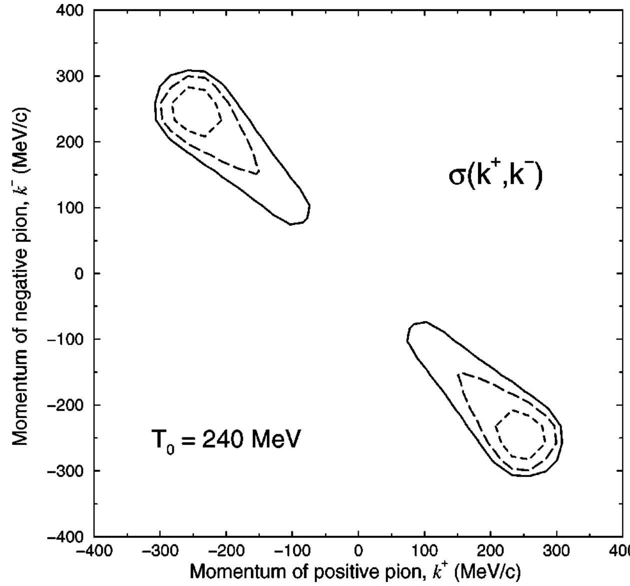

FIG. 2. Contours of the correlation coefficient for pion pairs of opposite charge as a function of the final momenta, for the one-dimensional scenario considered in Ref. [12] which uses a space and time dependent effective mass emulating the results obtained in Ref. [7] for the expansion of a system with the initial temperature of $T_0 = 240$ MeV and an initial radius of $R_0 = 5$ fm. The contours are separated by factors of 2.

It may be worthwhile to note that the visibility of the signal may be enhanced if the coincidence rate is divided by the product of the corresponding single-particle rates, $\langle \hat{N}_{\mathbf{k}'}^{\nu'} \hat{N}_{\mathbf{k}}^{\nu} \rangle / (\langle \hat{N}_{\mathbf{k}'}^{\nu'} \rangle \langle \hat{N}_{\mathbf{k}}^{\nu} \rangle)$. This procedure eliminates the overall variation due to the dependence of the singles yields on energy and angle, and the ratio unity in those kinematical regions where no correlations are present.

It is also important to emphasize that in each individual pion-pair emission process, it is entirely random whether a given partner (e.g., the positive one) goes forwards or backwards. Thus no overall charge separation occurs, but whenever one pion is emitted in one direction then its partner is emitted (approximately) oppositely.

In practice, when making an observation, it is not possible to identify a single quantum mode $\mathbf{k}$. Rather, a collection is made over a phase-space domain that encompasses many individual modes. In that situation, one may define the ap-

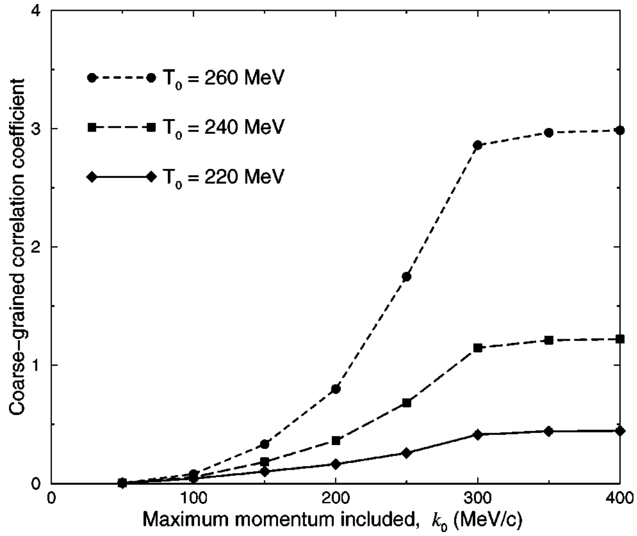

FIG. 3. The coarse-grained correlation coefficient for oppositely moving charge-conjugate soft pions as a function of the cutoff momentum k_0, for three different initial scenarios.

propriate coarse-grained particle number operator, $\hat{\mathcal{N}}^{\nu}_{\mathcal{K}}(t) = \Sigma_{\mathbf{k}\in\mathcal{K}}\hat{N}^{\nu}_{\mathbf{k}}(t)$, where $\mathcal{K}$ denotes a certain domain of $\mathbf{k}$ values. The associated coarse-grained correlation coefficient is then of the form

$$\Sigma^{\nu'\nu}_{\mathcal{K}'\mathcal{K}} \equiv \langle \hat{\mathcal{N}}^{\nu'}_{\mathcal{K}'}(t)\hat{\mathcal{N}}^{\nu}_{\mathcal{K}}(t)\rangle - \langle \hat{\mathcal{N}}^{\nu'}_{\mathcal{K}'}(t)\rangle\langle \hat{\mathcal{N}}^{\nu}_{\mathcal{K}}(t)\rangle. \tag{5}$$

To examine the effect of such a coarse graining, we let $\mathcal{K}$ encompass all positive pions moving forward with a momentum up to k_0 while $\mathcal{K}'$ includes all negative pions moving backward with a momentum down to $-k_0$.

Figure 3 shows the resulting coarse-grained correlation coefficient for three scenarios corresponding to different initial degrees of agitation and restoration. In all cases the coarse-grained correlation coefficient saturates once the cutoff k_0 is above ≈ 300 MeV/c, the maximum momentum for which the amplification mechanism is effective. Thus there is no gain in signal by extending the cutoff to higher values, while the background increases. In fact, in order to reduce contamination from ρ decays (see below), it may be desirable to employ a momentum cutoff below 300 MeV/c and the results suggest that this might well be possible. In the other direction, at the lower momenta, it is practically difficult to measure pions with $k_\perp$ much below 100 MeV/c but these are seen to contribute relatively little to the total signal.

Another important feature brought out in Fig. 3 is the good sensitivity of the signal to the physical scenario. For the relatively moderate range of initial temperatures considered, $T_0 = 220,240,260$ MeV, the initial chiral order parameter takes on the values $\sigma_0 = 19,26,42$ MeV and these differences have become magnified in the resulting correlation signal. This feature demonstrates that the signal is in fact quite sensitive to the degree of chiral restoration achieved early on.

The purpose of this Rapid Communication is to draw attention to the possible utility of analyzing the large-angle correlations between soft charge-conjugate pion pairs. The focus has been on arguing that one would expect an effect on rather general grounds, based simply on the nonequilibrium evolution of the chiral order parameter which in turn endows the pion with a temporal modulation of its effective mass that provides the possibility for parametric amplification. Moreover, within the simple linear σ model, we have tried to achieve an idea of the magnitude and persistence of the expected effect. Of course, this ‘‘signal’’ is partially obscured or eroded by a number of other processes and thus any attempt to extract it from the experimental data must take careful account of such ‘‘background’’ contributions.

One concern is the possible degradation of the primary signal in the course of the propagation of the pions as they leave the interaction zone. As the signal-carrying pions are soft they might easily be deflected significantly by encounters with other hadrons in the later stage of the reaction process. However, the very mechanism by which they are produced, the oscillatory relaxation of the chiral order parameter, guarantees that they appear only relatively late. Simple estimates, as well as more elaborate numerical studies [12], show that it takes several to many fm/c for the amplification mechanism to complete (say 5 fm/c). At this time, most of the material in the collision zone has already dispersed, so the DCC pions will be born in a relatively empty environment (which deviates from the regular vacuum principally by the deviation of the order parameter from its vacuum value). Moreover, since they are soft, they will have little chance of catching up with the expanding shell of regular collision debris. Thus, one may expect that the pions of interest will in fact propagate to the detector relatively undisturbed. These features are rather similar to those of the ‘‘Baked Alaska’’ scenario discussed by Bjorken [3]. In this connection it may be noted that any complications from elliptic flow are expected to be minimized, insofar as the analysis is being made for pions emitted in the mid-rapidity region of central collisions.

Another important issue concerns the possible presence of other agencies that may lead to a similar signal and thus obscure the signal of interest. While there are many physically different sources of charge-conjugate pion pairs, fortunately only few lead to strong back-to-back correlations. It is particularly important to discuss the dominant decay of $\rho(770), \rho \to \pi^+\pi^-$. Although the two pions are back-to-back correlated, they emerge from a ρ meson at rest with momenta of about 360 MeV/c which is somewhat above the upper limit of the expected effect ($k_{\max} \approx 300$ MeV/c). Moreover, while the decay of a ρ meson in motion may well contribute a single pion to the yield below this limit, the contribution to the coincidence yield is insignificant, even when taking account of the final ρ width and the in-medium thermal distortion. Thus, in the relevant domain of soft pions, the ρ decays are not expected to obscure to the two-body DCC signal.

It should also be noted that although $\eta(550)$ and $\omega(780)$ may contribute $\pi^+\pi^-$ pairs, these all arise in three-body decays which renders them only rather weakly correlated and so they should not pose a serious problem.

In conclusion, then, we suggest that the data now being taken at RHIC be analyzed for indications of the described signature in the large-angle correlation of soft charge-conjugate pion pairs. It may also be worthwhile to scrutinize existing SPS data for this signal. If indeed identified, this signal may offer a means for probing the degree of chiral restoration achieved and the subsequent DCC dynamics.

As a final comment, we wish to note that the considered effect should also manifest itself in electromagnetic observables, photons, and dileptons. For one thing, the effect is equally well present in the neutral pions which generally decay into a photon-pair before detection. Though it is impractical to reconstruct their predecay momenta, the excess of soft neutral pions should make a corresponding contribution to the photon yield. Moreover charge-conjugate pion pairs may annihilate to produce either photons or dileptons and the soft excess pions discussed here would then contribute accordingly to these observables. Therefore, should the analysis of the charged-pion data suggest the presence of the effect, this finding must be taken into account in the analysis of the electromagnetic signals.

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